

Exotic Decays of Strangelets¹

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Abstract

The stability of strange matter depends on the implicit assumption that baryon number is conserved. We examine the relevant effective operators which allow strangelets to decay by violating the conservation of baryon number. From the experimental lower limit of 10^{25} years on the stability of nuclei, we find a lower limit of the order 10^6 years on the stability of strangelets against such exotic decays.

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1 Introduction

The possibility in principle of absolutely stable matter containing strange quarks[1, 2] has generated a great deal of interest across many subfields of physics. A crucial implicit assumption for the stability of strange matter is that baryon number B is exactly conserved. Of course, we know nuclei are stable against $\Delta B \neq 0$ decays. The best mode-independent experimental lower limit to date is 1.6×10^{25} years.[3] However, there may be effective $\Delta B \neq 0$ interactions which are highly suppressed for nuclei but not for strangelets. In the following we examine such a hypothesis and show that whereas it is possible for strangelets to decay by reducing their baryon number, such a lifetime should be longer than 10^6 years.

In this report, we first review briefly the present experimental constraints on $\Delta B \neq 0$ interactions. We then point out the possible consequences of an effective $\Delta B = 2$, $\Delta N_s = 4$ operator on the stability of strange matter. We proceed to discuss some possible theoretical origins of such an operator, including R -parity nonconserving terms in supersymmetric extensions of the standard model. We conclude by deriving a phenomenological lower limit of 10^6 years on the lifetime of strangelets against these exotic decays.

2 Proton Decay

In order for the proton to decay, there has to be at least one fermion with a mass below that of the proton. Since only leptons have this property, the selection rule $\Delta B = 1$, $\Delta L = 1$ is applicable. The most studied decay mode experimentally is $p \rightarrow \pi^0 e^+$, which requires the effective interaction

$$\mathcal{H}_{int} \sim \frac{1}{M^2}(uude). \tag{1}$$

Note that because the above effective operator involves 4 fermions, it is of dimension 6, hence the interaction must depend on some large mass M to the power -2 . The fact that[3]

$$\tau_{exp}(p \rightarrow \pi^0 e^+) > 9 \times 10^{32} \text{y} \quad (2)$$

implies that $M > 10^{16}$ GeV. This means that proton decay probes physics at the grand-unification energy scale.

3 Neutron-Antineutron Oscillation and Related Processes

The next simplest class of effective interactions has the selection rule $\Delta B = 2$, $\Delta L = 0$. The most well-known example is of course

$$\mathcal{H}_{int} \sim \frac{1}{M^5} (udd)^2, \quad (3)$$

which induces neutron-antineutron oscillation and allows a nucleus to decay by the annihilation of two of its nucleons. Note that since 6 fermions are now involved, the effective operator is of dimension 9, hence M appears to the power -5 . This means that for the same level of nuclear stability, the lower bound on M will be very much less than 10^{16} GeV.

The present best experimental lower limit on the $n - \bar{n}$ oscillation lifetime is[4] $8.6 \times 10^7 \text{s}$. Direct search for $NN \rightarrow \pi\pi$ decay in iron yields[5] a limit of $6.8 \times 10^{30} \text{y}$. Although the above two numbers differ by 30 orders of magnitude, it is well established by general arguments as well as detailed nuclear model calculations that[6]

$$\tau(n\bar{n}) = 8.6 \times 10^7 \text{s} \Rightarrow \tau(NN \rightarrow \pi\pi) \sim 2 \times 10^{31} \text{y}. \quad (4)$$

Hence the two limits are comparable. To estimate the magnitude of M , we use

$$\tau(n\bar{n}) = M^5 |\psi(0)|^{-4}, \quad (5)$$

where the effective wavefunction at the origin is roughly given by[7]

$$|\psi(0)|^{-2} \sim \pi R^3, \quad R \sim 1 \text{ fm.} \quad (6)$$

Hence we obtain $M > 2.4 \times 10^5 \text{ GeV}$. Bearing in mind that \mathcal{H}_{int} is likely to be suppressed also by products of couplings less than unity, this means that the stability of nuclei is sensitive to new physics at an energy scale not too far above the electroweak scale of 10^2 GeV . It may thus have the hope of future direct experimental exploration.

4 Effective Interactions Involving Strangeness

Consider next the effective interaction

$$\mathcal{H}_{int} \sim \frac{1}{M^5} (uds)^2. \quad (7)$$

This has the selection rule $\Delta B = 2$, $\Delta N_s = 2$ ($\Delta S = -2$). In this case, a nucleus may decay by the process $NN \rightarrow KK$ [8]. Although such decay modes have never been observed, the mode-independent stability lifetime of $1.6 \times 10^{25} \text{ y}$ mentioned already is enough to guarantee that it is negligible. However, consider now

$$\mathcal{H}_{int} = \frac{1}{M^5} (uss)^2, \quad (8)$$

which has the selection rule

$$\Delta B = 2, \quad \Delta N_s = 4. \quad (9)$$

To get rid of four units of strangeness, the nucleus must now convert two nucleons into four kaons, but that is kinematically impossible. Hence the severe constraint from the stability of nuclei does not seem to apply here, and the following interesting possibility may occur.

Strangelets with atomic number A (which is of course the same as baryon number B) and number of strange quarks N_s may now decay into other strangelets with two less units

of A and one to three less s quarks. For example,

$$(A, N_s) \rightarrow (A - 2, N_s - 2) + KK, \quad (10)$$

$$(A, N_s) \rightarrow (A - 2, N_s - 1) + KKK, \text{ etc.} \quad (11)$$

For stable strangelets, model calculations show[9] that $N_s \sim 0.8A$, hence the above decay modes are efficient ways of reducing all would-be stable strangelets to those of the smallest A .

Unlike nuclei which are most stable for A near that of iron, the energy per baryon number of strangelets decreases with increasing A . This has led to the intriguing speculation that there are stable macroscopic lumps of strange matter in the Universe. On the other hand, it is not clear how such matter would form, because there are no stable building blocks such as hydrogen and helium which are essential for the formation of heavy nuclei. In any case, the above exotic interaction would allow strange matter to dissipate into smaller and smaller units, until A becomes too small for the strangelet itself to be stable. A sample calculation by Madsen[9] shows that

$$m_0(A) - m_0(A - 2) \simeq (1704 + 111A^{-1/3} + 161A^{-2/3}) \text{ MeV}, \quad (12)$$

assuming $m_s = 100$ MeV, and $B^{1/4} = 145$ MeV. Hence the KK and KKK decays would continue until the ground-state mass m_0 exceeds the condition that the energy per baryon number is less than 930 MeV, which happens at around $A = 13$.

5 Theoretical Prognosis

If the standard model of quarks and leptons is extended to include supersymmetry, the $\Delta B \neq 0$ terms $\lambda_{ijk} u_i^c d_j^c d_k^c$ are allowed in the superpotential. In the above notation, all chiral superfields are assumed to be left-handed, hence q^c denotes the left-handed charge-conjugated

quark, or equivalently the right-handed quark, and the subscripts refer to families, *i.e.* u_i for (u, c, t) and d_j for (d, s, b) . Since all quarks are color triplets and the interaction must be a singlet which is antisymmetric in color, the two d quark superfields must belong to different families. (In the minimal supersymmetric standard model, these terms are forbidden by the imposition of R -parity. However, this assumption is not mandatory and there is a vast body of recent literature exploring the consequences of R -parity nonconservation.)

To generate an effective $(uss)^2$ operator, we need to evaluate the one-loop diagram for

$$ss \rightarrow \tilde{b}\tilde{b}, \quad (13)$$

where \tilde{q} denotes the supersymmetric scalar partner of q , and attach the Yukawa interactions

$$\lambda_{usb} u^c s^c \tilde{b}^c \quad (14)$$

to the two \tilde{b} 's. This is analogous to a recent calculation[10] of the effective $(udd)^2$ operator and involves the exchange of the t quark, the W boson, and their supersymmetric partners \tilde{t} , and \tilde{w} . The end result is

$$\mathcal{H}_{int} = \frac{3g^4 \lambda_{usb}^2 (\tilde{A} m_b)^2 m_{\tilde{w}}}{8\pi^2 m_{\tilde{b}}^4 m_{\tilde{b}^c}^4} |V_{ts}|^2 F(m_t^2, m_{\tilde{t}}^2, m_W^2, m_{\tilde{w}}^2), \quad (15)$$

where the $\tilde{b}\tilde{b}^c$ mass term is assumed to be $\tilde{A} m_b$, V_{ts} is the quark-mixing matrix entry for t to s through the W boson, and F is a known function of the four internal masses in the loop. Using the correspondence of $n - \bar{n}$ oscillation to NN annihilation inside a nucleus, we estimate the lifetime of strangelets from the above effective interaction assuming $\tilde{A} = 200$ GeV, $\lambda_{usb} < 1$, and both scalar quark masses greater than 200 GeV, to be

$$\tau > 10^{20} \text{y}. \quad (16)$$

This tells us that such contributions are negligible from the supersymmetric standard model.

If we extend the supersymmetric standard model to include additional particles belonging to the fundamental **27** representation of E_6 , inspired by superstring theory, then it

is possible[11] to obtain an effective $(uss)^2$ operator without going through a loop. Two new interactions $us\tilde{h}$ and $\tilde{h}s^cN$ are now possible, where \tilde{h} is a new color-triplet scalar of charge $-1/3$ and N is a new neutral color-singlet fermion. They combine to form an effective $M^{-5}(uss)^2$ operator if N is allowed a Majorana mass, thereby breaking baryon-number conservation. For the process

$$(A, N_s) \rightarrow (A - 2, N_s - 2) + KK, \quad (17)$$

we estimate the lifetime to be[10]

$$\tau \sim \frac{32\pi m_N^2}{9\rho_N} \left[\frac{M}{\tilde{\Lambda}} \right]^{10} \sim 1.2 \times 10^{-28} \left[\frac{M}{\tilde{\Lambda}} \right]^{10} \text{ y}, \quad (18)$$

where $m_N \sim 1 \text{ GeV}$, $\rho_N \sim 0.25 \text{ fm}^{-3}$ is the nuclear density, and $\tilde{\Lambda} \sim 0.3 \text{ GeV}$ is the effective interaction energy scale corresponding to using Eq. (6). If we assume $M = 1 \text{ TeV}$, then

$$\tau \sim 2 \times 10^7 \text{ y}. \quad (19)$$

6 Lower Limit on the Exotic Decay Lifetime of Strangelets

Since τ depends on $M/\tilde{\Lambda}$ to the power 10 in Eq. (18), it appears that a much shorter lifetime than that of Eq. (19) is theoretically possible. However, there is a crucial phenomenological constraint which we have yet to consider. Although two nucleons cannot annihilate inside a nucleus to produce four kaons, they can make three kaons plus a pion. The effective $(uss)^2$ operator must now be supplemented by a weak transition

$$s \rightarrow u + d + \bar{u}. \quad (20)$$

We can compare the effect of this on $NN \rightarrow KKK\pi$ versus that of the $(uds)^2$ operator on $NN \rightarrow KK$ discussed in Ref.[10]. We estimate the suppression factor to be

$$\frac{3 \times 10^{-3}}{(4\pi^2)^2} |V_{us}|^2 \tilde{\Lambda}^4 G_F^2 \sim 10^{-19}, \quad (21)$$

where G_F is the Fermi weak coupling constant. Since the stability of nuclei is at least 1.6×10^{25} years, this gives a firm lower limit

$$\tau > 10^6 \text{y} \tag{22}$$

on the lifetime of strangelets against $\Delta B = 2$, $\Delta N_s = 4$ decays.

7 Conclusion and Outlook

We have pointed out in this report that an effective $(uss)^2$ interaction may cause stable strange matter to decay, but the lifetime has a lower limit of 10^6 years. This result has no bearing on whether stable or metastable strangelets can be created and observed in the laboratory, but may be important for understanding whether there is stable strange bulk matter left in the Universe after the Big Bang and how it should be searched for. For example, instead of the usual radioactivity of unstable nuclei, strange matter may be long-lived kaon and pion emitters. Furthermore, if the particles mediating this effective interaction have masses of order 1 TeV as discussed, then forthcoming future high-energy accelerators such as the LHC at CERN will have a chance of confirming or refuting their existence.

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